Local Geometry of the Fermi Surface and Magnetoacoustic Responce of Two-Dimensional Electron Systems in Strong Magnetic Fields

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(Dated: February 2, 2008)

A semiclassical theory for magnetotrasport in a quantum Hall system near filling factor $\nu = 1/2$ based on the Composite Fermions physical picture is used to analyze the effect of local flattening of the Composite Fermion Fermi surface (CF-FS) upon magnetoacoustic oscllations. We report on calculations of the velocity shift and attenuation of a surface acoustic wave (SAW) which travels above the two-dimensional electron system, and we show that local geometry of the CF-FS could give rise to noticeable changes in the magnitude and phase of the oscillations. We predict these changes to be revealed in experiments, and to be used in further studies of the shape and symmetries of the CF-FS. Main conclusions reported here could be applied to analyze magnetotransport in quantum Hall systems at higher filling factors $\nu = 3/2, 5/2$ provided the Fermi-liquid-like state of the system.

PACS numbers: 71.10.pm, 72.55.+s, 73.43.-f

INTRODUCTION AND BACKGROUND

A two-dimensional electron gas (2DEG) in a strong magnetic field reveals rich and complex physics. Near half filling of the lowest Landau level ($\nu = 1/2$) the ground state of such a system is shown to be a compressible Fermi-liquid-like state of Composite Fermions (CF) [1]. These quasiparticles are distributed inside the Composite Fermion Fermi Surface (CF-FS). A similar physical picture could be adopted to describe the 2DEG at half filling of the next Landau level ($\nu = 3/2$). Experimental evidences of the CF Fermi sea at $\nu = 1/2, 3/2$ were repeatedly obtained during the recent decade [2].

For higher filling factors close to $\nu = N + 1/2$ where $N \geq 3$ is an integer, the exchange interaction would lead to an instability towards charge density wave formation in the relevant Landau levels. The ground state of the 2DEG for these filling factors corresponds to a charge density wave (CDW) and has a striped structure [3, 4, 5, 6]. It could be described as a sequence of one dimensional stripes alternating between the adjacent filling factors N and N + 1. This gives rise to a strikingly anisotropic transport properties of the 2DEG at half filling of higher Landau levels [7, 8] which were revealed in experiments [9, 10, 11].

The quantum Hall state at $\nu = 5/2$ is perhaps the most enigmatic due to its position in the magnetic field spectrum between the high Landau level $N \geq 3$ stripe phases and the low Landau level ($N \leq 1$) Fermi-liquid like states. Theoretical studies of this state started from the model of paired CFs [12]. Depending on the interaction strength within the system, the 2DEG at $\nu = 5/2$ could reveal a striped state, a Fermi–liquid, or the paired state [13]. Numerical simulations presented in [13] give grounds to believe that at $\nu = 5/2$ the CF Fermi liquid undergoes condensation to the paired state at low temperature limit. Also, there could be a transition from the Fermi–liquid to the striped phase when in-plane magnetic field is applied. Recently, an experimental evidence of the CF-FS at $\nu = 5/2$ was obtained [14].

So, both theory and experiment give grounds to believe that physical picture of CFs which form a Fermi sea could be succesfully employed to describe the 2DEG states at half filling of the lowest Landau levels ($N \leq 2$). However, the geometry of the CF-FS was not analyzed up to present. It is usually assumed that the CF Fermi liquid is isotropic, and the CF-FS is a circle in the two-dimensional quasimomenta space. This is an obvious oversimplification. Real samples commonly used in studies of the quantum Hall effect have 2DEGs deposited in GaAs/AlGaAs heterostructures. Therefore the crystalline field of the host semiconductor could significantly influence the CF-FS geometry distorting the original Fermi circle [15]. Another sourse of the CF-FS anisotropy, especially for higher filling factors ($\nu = 3/2, 5/2$), is the interaction in the electron system. The development of highly anisotropic charge-density wave formations (striped phases) at high filling factors including 5/2, gives us strong arguments to expect these interactions to work as an extra crystalline field at the Fermi-liquid state of the 2DEG at $\nu = 3/2, 5/2$. As a result the CF-FS shape could be further modified.

Theory of magnetotransport in metals showes that the FS local geometry noticeably affects the electron response of the metal to an external perturbation [16]. The change in the response occurs under the nonlocal regime of propagation of the disturbance when the mean free path of electrons l is large compared to the wavelength of the disturbance λ. The reason is that in this nonlocal regime only those electrons whose motion is somehow consistent with the propagating perturbation can strongly absorb its energy. These "efficient" electrons are concentrated on small "effective" segments of the FS.

When the FS includes flattened segments it leads to

an enhancement of the contribution from these segments to the electron density of states (DOS) on the FS. Usually this enhanced contribution is small compared to the main term of the DOS which originates from all the remaining parts of the FS. Therefore it cannot produce noticeable changes in the response of the metal under the local regime of propagation of the disturbance ($l << \lambda$) when all segments of the FS contribute to the response functions essentially equally. However the contribution to the DOS from the flattened section can be congruent to the contribution of a small "effective" segment of the FS. In other words when the curvature of the FS becomes zero at some points on an "effective" part of the FS it can give a significant enhancement of efficient electrons and, in consequence, a noticeable change in the response of the metal to the disturbance.

Due to the same reasons we can expect local geometrical features of the CF-FS to give significant effects on the 2DEG response to an external disturbance. As well as for convenient metals, these effects are to be revealed within a nonlocal regime ($l > \lambda$). It was shown before that the local flattening of the CF-FS could give rise to a strong anisotropy in the response of a 2DEG to a surface acoustic wave (SAW) [17]. Such anomaly was observed in experiments on modulated 2DEG near $\nu = 1/2$ [18].

Here, we analyze the influence of the CF-FS local geometry on so called geometric resonances which were repeatedly observed in 2DEGs in strong magnetic field [2], as well as in convenient metals [19]. These oscillations could be revealed within a nonlocal regime, and they appear due to periodical reproduction of the most favorable conditions for the resonance absorption of the energy of the external disturbance by quasiparticles at stationary points on the cyclotron orbit where they move along the wave front of the disturbance. When the external disturbance is associated with an acoustic wave these geometric resonances are also called magnetoacoustic oscillations [19]. In the following analysis we mostly consider magnetoacoustic oscillations in the 2DEG at $\nu = 1/2$ state, and we describe this state within the framework of Halperin, Lee and Read (HLR) theory [1]. However, we believe that the main results of the present analysis could be applied to study magnetotransport in 2DEGs at higher filling factors $(3/2, 5/2)$ provided that the system is at a Fermi-liquid-like state.

MAIN EQUATIONS

Due to the piezoelectric properties of GaAS, the velocity shift $(\Delta s/s)$ and the attenuation rate (Γ) for the SAW propagating along the x axis across the surface of a heterostructure containing 2DEG, take the form [20]:

$$\frac{\Delta s}{s} = \frac{\alpha^2}{2} \mathrm{Re} \left(1 + \frac{i\sigma_{xx}}{\sigma_m} \right)^{-1}; \tag{1}$$

$$\Gamma = -q \frac{\alpha^2}{2} \mathrm{Im} \left(1 + \frac{i\sigma_{xx}}{\sigma_m} \right)^{-1}. \tag{2}$$

Here $\mathbf{q}, \omega = sq$ are the SAW wave vector and frequency, respectively, α is the piezoelectric coupling constant, $\sigma_m = \varepsilon s/2\pi$, ε is an effective dielectric constant of the background and σ_{xx} is the component of the electron conductivity tensor.

According to HLR theory, the electron resistivity tensor ρ at $\nu = 1/2$ is given by:

$$\rho = \sigma^{-1} = \rho^{CF} + \rho^{CS} \tag{3}$$

where ρ^{CF} is the CF resistivity tensor, and the contribution ρ^{CS} originates from the Chern–Simons formulation of the theory. This tensor contains only off diagonal elements $\rho_{xy}^{CS} = -\rho_{yx}^{CS} = 4\pi\hbar/e^2$.

The CF resistivity tensor ρ^{CF} could be calculated as the inverse for the CF conductivity $\tilde{\sigma}$ $(\rho^{CF} = \tilde{\sigma}^{-1})$. We carry out our analysis in a regime where $\rho_{xx}\rho_{yy} << \rho_{xy}^2$, therefore the relevant component of the electron conductivity could be written in the form:

$$\sigma_{xx}(q) = \frac{e^4}{(4\pi\hbar)^2} \rho_{yy}^{CF} = \frac{e^4}{(4\pi\hbar)^2} \frac{\tilde{\sigma}_{xx}}{\tilde{\sigma}_{xx}\tilde{\sigma}_{yy} + (\tilde{\sigma}_{xy})^2}. \tag{4}$$

The CFs are supposed to experience not actual but reduced magnetic field $B_{eff} = B - B_{1/2}$ where $B_{1/2}$ corresponds to one half filling of the lowest Landau level. Their motion could be described within a semiclassical approximation based on the Boltzmann's transport equation. Following standard methods [21] we obtain:

$$\tilde{\sigma}_{\alpha\beta} = \frac{m^* e^2}{2\pi\hbar^2} \sum_n \frac{v_{n\beta}(-q) v_{n\alpha}(q)}{in\Omega - i\omega + 1/\tau}. \tag{5}$$

Here, m^*, Ω are the CF cyclotron mass and their cyclotron frequency at the field B_{eff}; τ is the CF scattering time, and $v_{n\alpha}(q)$ are the Fourier transforms of the CF velocity components:

$$v_{nx}(q) = \frac{n}{2\pi} \frac{\Omega}{q} \int_0^{2\pi} d\psi$$

$$\times \exp\left\{ in\psi - \frac{iq}{\Omega} \int_0^{\psi} v_x(\psi') d\psi' \right\}; \tag{6}$$

$$v_{ny}(q) = \frac{1}{2\pi}\int_0^{2\pi} d\psi v_y(\psi)$$

$$\times \exp\left\{ in\psi - \frac{iq}{\Omega}\int_0^{\psi} v_x(\psi')d\psi' \right\}. \qquad (7)$$

The variable ψ included in these expressions is the angular coordinate of the CF cyclotron orbit.

The most favorable conditions for magnetoacoustic oscillations to be revealed occur at moderately strong effective magnetic field when $ql >> \Omega\tau >> 1$. Under these conditions the main contributions to the integrals over ψ in the expressions (6),(7) come from the neighborhoods of stationary points at CF cyclotron orbits. So, the expressions for $v_{n\alpha}(q)$ could be rewritten as:

$$v_{nx}(q) = \frac{n\Omega}{q}\cos\left(qR - \frac{\pi n}{2} - \Phi\right) X(qR); \qquad (8)$$

$$v_{ny}(q) = -i\sin\left(qR - \frac{\pi n}{2} - \Phi\right) V(qR). \qquad (9)$$

Here, $2R$ is the diameter of the CF cyclotron orbit in the direction of propagation of the SAW, $X(qR)$ is a dimensionless quantity, and $V(qR)$ has dimensions of velocity. Both $X(qR)$ and $V(qR)$ are power functions of qR with the same exponent. As we show below, the value of the exponent is determined with the local geometry of small effective segments of the CF-FS which correspond to the vicinities of the stationary points. When these segments are flattened, this leads to significant changes in the magnitude of the magnetoacoustic oscillations.

THE CF-FS MODEL

Within the commonly used isotropic model of CF Fermi liquid at $\nu = 1/2$ the CF-FS is a circle, and its radius p_F equals $\sqrt{4\pi N\hbar^2}$ where "N" is the electron density. To develop more realistic model of the CF-FS we include a periodic static electric field applied along the "y" direction which provides CFs with the potential energy of magnitude U_g ($\mathbf{g}$ is the wave vector of the electric field). The above electric field could originate from interactions with electrons of lower Landau levels at $\nu = 3/2, 5/2$ and from the crystalline field of the host semiconductor. The latter is especially important at $\nu = 1/2$. The point is that wherever it comes from, this field distorts the CF-FS, including formation of local anomalies of the CF-FS curvature.

Assume for simplicity that the electric modulation is weak ($U_g << E_F$, where E_F is the CFs Fermi energy). Then we can use the nearly-free-electron model to derive

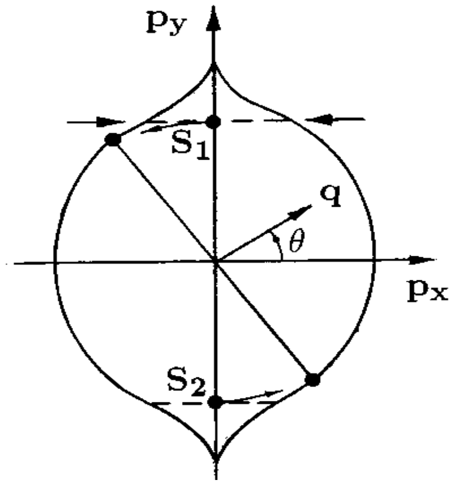

FIG. 1: The shape of the CF-FS in the nearly-free-electron approximation (solid line), and the CF-FS described with the Eq. (14) (dashed line). Point S_1 and S_2 are associated with the stationary points at the CF cyclotron orbit when the SAW wave vector $\mathbf{q}$ points in the p_x direction.

the energy-momentum relation for the CFs. When the modulation period is small enough $\hbar g > 2p_F$ we obtain:

$$E(\mathbf{p}) = \frac{p_x^2}{2m} + \frac{p_y^{*2}}{2m} + \frac{(\hbar g)^2}{8m} - \sqrt{\left(\frac{\hbar g p_y^*}{2m}\right)^2 + U_g^2}. \qquad (10)$$

Here $p_y^* = p_y - \hbar g/2$, m is the CF effective mass; U_g is the magnitude of the quasiparticle potential energy in the periodic electric field. Calculating the FS curvature:

$$K = \frac{1}{v^3}\left(2v_x v_y \frac{\partial v_x}{\partial p_y} - v_x^2\frac{\partial v_y}{\partial p_y} - v_y^2\frac{\partial v_x}{\partial p_x}\right) \qquad (11)$$

with $v = \sqrt{v_x^2 + v_y^2}$ one can find it tending to zero when p_x tends to $\pm p_F(U_g/E_F)^{1/2}$ (See Fig.1).

In the vicinities of the corresponding points on the FS the quasiparticles velocities are nearly parallel to the the y direction. Near these zero curvature points we can derive an asymptotic expression for energy-momentum relation (10). Introducing (p_{x_0}, p_{y_0}) by $p_{x_0} = \zeta p_F$, $p_{y_0} = p_F\left(1 - \frac{1}{\sqrt{2}}\zeta^2\right)$, where $\zeta = \sqrt{U_g/E_F}$, we can expand the variable p_y in powers of $(p_x - p_{x_0})$, and keep the lowest order terms in the expansion. We obtain:

$$p_y - p_{y_0} = -\zeta(p_x - p_{x_0}) - \frac{2}{\zeta^4}\frac{(p_x - p_{x_0})^3}{p_F^2}. \qquad (12)$$

Near p_{x0}, where ($|p_x - p_{x_0}| < \zeta^2 p_F$) the first term on the right side of Eq.(12) is small compared to the second one and can be omitted. So we have:

$$E(\mathbf{p}) = \frac{4}{\zeta^4}\frac{p_F^2}{2m}\left(\frac{p_x - p_{x_0}}{p_F}\right)^3 + \frac{p_y^2}{2m}. \qquad (13)$$

The "nearly free" particle model can be used when ζ^2 is very small. For larger U_g the local flattening of the CF-FS can be more significant. To analyze the contribution to the conductivity from these flattened parts we generalize the expression for $E(\mathbf{p})$ and define our dispersion as:

$$E(\mathbf{p}) = \frac{p_0^2}{2m_1}\left|\frac{p_x}{p_0}\right|^\gamma + \frac{p_y^2}{2m_2}, \qquad (14)$$

where p_0 is a constant with the dimensions of momentum, m_i are the effective masses, and γ is a dimensionless parameter which determines the shape of the CF–FS . When $\gamma > 2$ the CF–FS looks like an ellipse flattened near the vertices $(0, \pm\sqrt{m_2/m_1}p_0)$. Near these points the curvature is:

$$K = -\frac{\gamma(\gamma-1)}{2p_0\sqrt{m_1/m_2}}\left|\frac{p_x}{p_0}\right|^{\gamma-2} \qquad (15)$$

and, $K \to 0$ at $p_x \to 0$. The CF-FS will be the flatter at $p_x = 0$, the larger is the parameter γ. A separate investigation is required to establish how γ depends on U_g. Here we postulate Eq.(14) as a natural generalization of Eq.(13).

When $p_F > \hbar g$ we have to consider the CF–FS as consisting of several branches belonging to several "bands" or Brillouin zones. The modulating potential wave vector $\mathbf{g}$ in this case determines the size of the "unit cell". However with this condition we also may expect some branches of the CF–FS to be flattened. Within an appropriate geometry of an experiment these flattened segments of the CF-FS become the effective parts of the FS. Consequently, the response of the CF system to the SAW could undergo significant changes. Prior to start the analysis of these changes we remark that our model of the deformed CF-FS (14) could be easily generalized and accomodated to more complicated geometry of electric field which determines the CF-FS shape and symmetries. However, even the simple model (14) captures the essential physics, enabling to include local flattenings of the CF-FS into consideration. Therefore we adopt this model in the further analysis.

RESULTS AND DISCUSSION

When the SAW propagates along the "x" direction the vertices S_1, S_2 of the flattened ellipse (14) correspond to the stationary points on the CF cyclotron orbits (see Fig. 1). The enhanced DOS of quasiparticles in their vicinities influences the features of the magnetoacoustic oscillations. Using the stationary-phase method we obtain the following asymptotics for the functions $X(qR)$ and $V(qR)$:

$$X(qR) = \frac{m^*}{p_0}V(qR)$$

$$= \frac{2}{\pi}\left(\frac{m^*}{\sqrt{m_1 m_2}}\right)^{1/\gamma}\frac{\Gamma(1/\gamma)}{\gamma}\left(\frac{2}{qR}\right)^{1/\gamma}. \qquad (16)$$

Here, $R = \dfrac{cp_0}{eB_{eff}}\sqrt{\dfrac{m_2}{m_1}}$, $\Gamma(x)$ is the gamma function. As for the phase shift Φ, it appears to be equal $\pi/2\gamma$.

Using these results (16), as well as standard formulas [22]:

$$\sum_{n=-\infty}^{\infty}\frac{1}{\omega + i/\tau - l\Omega} = -\frac{i\pi}{\Omega}\coth\frac{\pi(1-i\omega\tau)}{\Omega\tau}; \qquad (17)$$

$$\sum_{n=-\infty}^{\infty}\frac{(-1)^n}{\omega + i/\tau - l\Omega}$$

$$= -\frac{i\pi}{\Omega}\frac{1}{\sinh\left[\pi(1-i\omega\tau)/\Omega\tau\right]} \qquad (18)$$

we can transform the expressions (5) for the CF conductivity components to the form:

$$\tilde{\sigma}_{xx} = \frac{2}{\rho_0}\frac{b^2(1-i\omega\tau)}{(ql)^2}; \qquad (19)$$

$$\tilde{\sigma}_{xy} = -\sigma_{yx} = -\frac{2}{\rho_0}\frac{g^2}{(ql)^2}\left(\frac{qR}{2}\right)^{1-2/\gamma}$$

$$\times\frac{(1-i\omega\tau)\sin\left(2qR - \pi/\gamma\right)}{\sinh\left[\pi(1-i\omega\tau)/\Omega\tau\right]}; \qquad (20)$$

$$\tilde{\sigma}_{yy} = \frac{2}{\rho_0}\frac{d^2}{ql}\left(\frac{qR}{2}\right)^{1-2/\gamma}\left\{\coth\frac{\pi(1-i\omega\tau)}{\Omega\tau}\right.$$

$$\left. - \cos\left(2qR - \frac{\pi}{\gamma}\right)\sinh^{-1}\frac{\pi(1-i\omega\tau)}{\Omega\tau}\right\}. \qquad (21)$$

where $\rho_0 = m^*/Ne^2\tau$ is the CFs Drude resistivity; $l = \dfrac{\tau}{m^*}\sqrt{\dfrac{A}{\pi}}$; and A is the area of the CF-FS. The factors

b^2, d^2 and g^2 included in (20)–(22) are the dimensionless constants of the order of unity. Expressions for these constants are omitted for brevity. For a circular CF-Fs we have: $\gamma = 2,\ p_0 = p_F,\ b^2 = g^2 = d^2 = 1$, and our expressions (19)–(21) take on the form:

$$\tilde{\sigma}_{xx} = \frac{2}{\rho_0}\frac{1 - i\omega\tau}{(ql)^2}; \tag{22}$$

$$\tilde{\sigma}_{xy} = -\tilde{\sigma}_{yx} = \frac{2}{\rho_0}\frac{1 - i\omega\tau}{(ql)^2}\cos(2qR)$$

$$\times \sinh^{-1}\left[\frac{\pi(1 - i\omega\tau)}{\Omega\tau}\right]; \tag{23}$$

$$\tilde{\sigma}_{yy} = \frac{2}{\rho_0}\frac{1}{ql}\left\{\coth\left[\frac{\pi(1 - i\omega\tau)}{\Omega\tau}\right]\right.$$

$$\left. - \sin(2qR)\sinh^{-1}\left[\frac{\pi(1 - i\omega\tau)}{\Omega\tau}\right]\right\}. \tag{24}$$

Comparison of our expressions (19)–(21) with the results for a Fermi circle showes that the local flattening of the CF-FS near the points which correspond to the stationary points of the CFs cyclotron orbit enhances the magnitude of the magnetoacoustic oscillations of the CF conductivities. A similar effect was studied before for conventional metals [23]. The effect originates from the enhancement of the quasiparticle DOS at flattened segments of the FS.

The above considered enhancement of magnetoacoustic geometric oscillations could be manifested only when the stationary points on the CF cyclotron orbit correspond to the points located at flattened segments of the CF-FS. Therefore the effect has to be very sensitive to variations in the direction of the SAW propagation. Suppose that the SAW travels at some angle θ with respect to the symmetry axis of the CF-FS as it is shown in Fig. 1. Then the stationary points slip from the flattened pieces and fall into "normal" segments of the FS whose curvature takes on nonzero values. Due to the lower DOS of quasiparticles at these "normal" CF-FS segments, the number of efficient CFs which can participate in the absoption of the SAW energy decreases when the angle θ increases. This results in the noticeable reduction of the oscillations.

Assuming a nonzero value for the angle θ, we can present Fourier transforms of the CF velocity components in the form:

$$v_{nx}(q) = \frac{n\Omega}{q}\left[\cos\left(qR - \frac{\pi n}{2}\right)S_\gamma(qR,\theta)\right.$$

$$\left. + \sin\left(qR - \frac{\pi n}{2}\right)W_\gamma(qR,\theta)\right]; \tag{25}$$

$$v_{ny}(q) = -\frac{ip_0}{m^*}\left[\sin\left(qR - \frac{\pi n}{2}\right)S_\gamma(qR,\theta)\right.$$

$$\left. - \cos\left(qR - \frac{\pi n}{2}\right)W_\gamma(qR,\theta)\right]. \tag{26}$$

Here,

$$S_\gamma(qR,\theta) = \frac{2}{\pi}\left(\frac{m^*}{\sqrt{m_1 m_2}}\right)^{1/\gamma}$$

$$\times \int_0^\infty \cos\left[\frac{qR}{2}\left(\frac{m_2}{m_1}\sin^2\theta y^2 + \cos^\gamma\theta y^\gamma\right)\right]; \tag{27}$$

$$W_\gamma(qR,\theta) = \frac{2}{\pi}\left(\frac{m^*}{\sqrt{m_1 m_2}}\right)^{1/\gamma}$$

$$\times \int_0^\infty \sin\left[\frac{qR}{2}\left(\frac{m_2}{m_1}\sin^2\theta y^2 + \cos^\gamma\theta y^\gamma\right)\right]. \tag{28}$$

Now, we expaind these functions S_γ in series in powers of a dimensionless parameter ξ $\left(\xi = \frac{m_2}{m_1}\left(\frac{qR}{2}\right)^{1-2/\gamma} \times \tan^2\theta\right)$. For small angles θ we have $\xi << 1$, and the expansion accepts the form [22]:

$$S_\gamma(qR,\theta) = \frac{1}{\gamma\cos\theta}\left(\frac{2}{qR}\right)^{1/\gamma}$$

$$\times \sum_{r=0}^{\infty}\frac{(-1)^r}{r!}\xi^r\Gamma\left(\frac{2r+1}{\gamma}\right)\cos\left[\pi\frac{1 - r(\gamma-2)}{2\gamma}\right]. \tag{29}$$

As θ increases to the values guaranteeing the inequality $\xi > 1$ to be valid, we have to use a different power expansion, namely:

$$S_\gamma(qR,\theta) = \frac{1}{2\sin\theta}\sqrt{\frac{m_1}{m_2}\frac{2}{qR}}\sum_{r=0}^{\infty}\frac{(-1)^r}{r!}\xi^{-\gamma r/2}$$

$$\times\Gamma\left(\frac{\gamma r + 1}{2}\right)\cos\left[\frac{\pi}{4}\left(r(\gamma-2) + 1\right)\right]. \tag{30}$$

Expansions for the function $W_\gamma(qR,\theta)$ could be ontained from (29),(30) by replacing cosines by sines of the same angles.

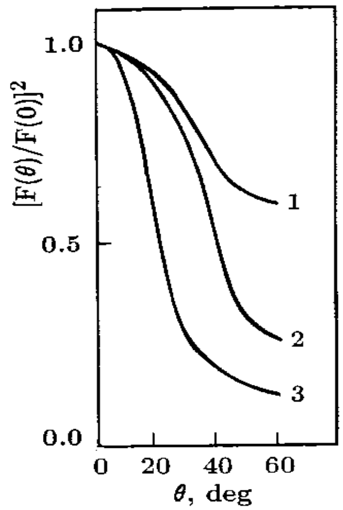

FIG. 2: Angular dependence of the amplitude of magnetoacoustic oscillations. The curves are plotted for $qR = 10$; $m_1 = m_2$; $\gamma = 4$ (curve 1); $\gamma = 6$ (curve 2); and $\gamma = 8$ (curve 3).

When $\theta \to 0$ we arrive at our former asymptotic expressions for the CF velocity components (16). While θ increases, the functions $S_\gamma(qR,\theta)$ and $W_\gamma(qR,\theta)$ diminish and approach the value $(\pi m_1/4m_2 qR)^{1/2}$ which is a typical estimate for an elliptical CF-FS as $\theta \to 90^o$. Angular dependence of the amplitude factor of magnetoacoustic oscillations $F_\gamma^2(qR,\theta) = S_\gamma^2(qR,\theta) + W_\gamma^2(qR,\theta)$ is presented in the Fig. 2. We see in this figure that the effect of local flattening of the CF-FS on the oscillations amplitude remains distinguishable even for moderate flattenings ($\gamma = 4$).

Differences in magnitudes of geometric osclillations of the CF conductivity components are manifested in the electron conductivity. Substituting our results for $\tilde{\sigma}_{\alpha\beta}$ into (4) we have:

$$\sigma_{xx} = \frac{qe^2}{8d^2 p_0}\left(\frac{2}{qR}\right)^{1-2/\gamma}$$

$$\times \frac{\sinh[\pi(1-i\omega\tau)/\Omega\tau]}{\cosh[\pi(1-i\omega\tau)/\Omega\tau] - \cos(2qR - \pi/\gamma)}. \tag{31}$$

We compare this expression with the corresponding result for a circular CF-FS

$$\sigma_{xx}^{el} = \frac{qe^2}{8p_F}\frac{\sinh[\pi(1-i\omega\tau)/\Omega\tau]}{\cosh[\pi(1-i\omega\tau)/\Omega\tau] - \sin(2qR)}, \tag{32}$$

and we see that both amplitude and phase of the geometric oscillations in the electron conductivivty differ from those for the CF Fermi circle. The same could be applied to magnetoacoustic oscillations described with the expressions (1),(2). When the effective parts of the CF-FS are flattened, the amplitude of the magnetoacoustic oscillations drops. We also conclude that varying the direction of propagation of the SAW we can observe angular dependence of the oscillations amplitude. The latter originates from the angular dependence of the CF conductivities discussed before. We have grounds to expect it to be revealed in experiments.

Finally, we belive that Fermi-liquid state of a 2DEG in quantum Hall regime at $\nu = 1/2, 3/2, 5/2$ is anisotropic and exhibits an anisotropic CF-FS. The CF-FS geometry reflects symmetries of the crystalline fields of the host semiconductor. For higher filling factors $\nu = 3/2, 5/2$ it also could show effects of interactions in the electron system. It was already found that screening due to polarization of remote Landau levels plays an essential role for the preferred orientation of the stripes induced by an in-plane magnetic field at $\nu = 5/2$ [24]. Therefore, we may conjecture that at weaker in-plane fields when Fermi-liquid like state of the 2DEG still exits, these polarization effects could give extra anisotropies to the Cf-FS.

Anisotropic CF-FSs usually include some flattened segments. Even a naive model (10) based on the nearly-free-electron approach, demonstrates that local flattening of the CF-FS appears as a result of a weak deformation of the latter with an electric modulation. In general, local flattenings are initiated with electric fields acting within a 2DEG like crystalline fields in usual metals. Accordingly, locations of the flattened segments conform with the symmetries of the CF-FS and could reveal these symmetries. The results of the present analysis show that magnetoacoustic oscillations in the velocity shift and attenuation of the SAW travelling in piezoelectric GaAs/AlGaAs heterostructures above the 2DEG, could be used as a tool to discover local flattenings at the CF-FS when the 2DEG is in the quantum Hall regime at Fermi-liquid like state. This could give a new knowledge of the shape and symmetries of the CF-FS and, consequently, a better understanding of the magnetotransport in quantum Hall systems near half filling of lowest Landau levels. It would be especially interesting to compare symmetries of the CF-FS at the Fermi-liquid like state of the 2DEG at $\nu = 5/2$ with the characteristic symmetries of the striped state of the system at the same filling factor. It is possible that such comparison would give some unusual results providing a new insight in the nature of transition from the Fermi-liquid to the striped phase of the 2D electron system.

Acknowledgments: The author thanks G.M. Zimbovsky for help with the manuscript.

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